

Supersymmetric Signatures at an $e\gamma$ Collider

Ken Kiers^b, John N. Ng^{a,b} and Guohong Wu^a

^a*TRIUMF Theory Group*

4004 Wesbrook Mall, Vancouver, B.C., Canada V6T 2A3

and

^b*Department of Physics, University of British Columbia*

Vancouver, B.C., Canada V6T 1Z1

Abstract

High energy electron-photon colliders provide unique opportunities for probing physics beyond the standard model. We have studied the experimental signatures for two supersymmetric scenarios, with the lightest supersymmetric particle (LSP) being either the lightest neutralino or the gravitino. In the “neutralino LSP” scenario favored by the minimal supersymmetric standard model (MSSM), it is found that some basic parameters of the model, μ , $\tan\beta$, M_1 and M_2 , may be uniquely determined from the outgoing electron energy spectrum without assuming high scale unification of the masses or couplings. In the “gravitino LSP” scenario which occurs naturally in models of low energy dynamical supersymmetry breaking, it is possible to have background-free signatures if the next-to-lightest supersymmetric particle (NLSP) has a long decay length. In cases that the NLSP decays quickly, ways to distinguish among the experimental signatures of the two scenarios and of the standard model (SM) background are discussed.

I. INTRODUCTION

It was first pointed out by Ginzburg, *et al* [1] that an e^+e^- collider could be adapted to make an $e\gamma$ collider by back-scattering laser light off the high energy electron beam. The resulting $e\gamma$ collider would be expected to have a center of mass energy and luminosity comparable to that of the e^+e^- collider, and this would open up new avenues [2] in testing the standard model (SM) of particle physics and in discovering physics beyond the standard model.

Experimental signatures of supersymmetry [3,4] at high energy $e\gamma$ colliders have been investigated to some extent [5,6]. Assuming R -parity conservation, the lightest supersymmetric particle (LSP) is a stable neutral particle and will escape detection. The studies to date have focused on the simplest case of lightest neutralino (χ_1) production under the assumption that χ_1 is the LSP and is gaugino-like. The process is $e\gamma \rightarrow \tilde{e}\chi_1$ with subsequent decay of the selectron $\tilde{e} \rightarrow e\chi_1$, and the experimental signature is $e + \cancel{E}$. A well-known feature of this process is that it has a lower threshold for selectron production than the analogous e^+e^- process.

In this letter we extend the previous supersymmetric (SUSY) studies with $e\gamma$ colliders in two respects. Firstly, we will consider two scenarios, allowing the LSP to be either the lightest neutralino χ_1 or the gravitino. The “neutralino LSP” scenario is often assumed to be realized in the minimal supersymmetric standard model (MSSM) [4], whereas in models of low energy dynamical supersymmetry breaking [7] the LSP is the gravitino and the next-to-lightest supersymmetric particle (NLSP) may be the lightest neutralino or a right-handed slepton, whose decay proceeds through its coupling to the Goldstino component of the gravitino [8]. Some of the phenomenology of a light gravitino has recently been discussed in the context of e^+e^- and hadron colliders [9,10]. Secondly, we will generalize the previous analysis [5,6] in the MSSM (without assuming high scale unification of masses or couplings) by allowing for more production channels and decay modes of the \tilde{e} . We find that a com-

plete determination of those MSSM parameters that enter the neutralino and chargino mass matrices, μ , $\tan\beta$, M_1 , and M_2 (as well as the selectron mass), may be possible for certain ranges of these parameters by simply measuring the outgoing electron energy spectrum.

It is well known that the cross section for $e\gamma \rightarrow W\nu \rightarrow e\nu\bar{\nu}$ approaches a constant at asymptotic energies [11]; hence it is crucial to reduce the W background as much as possible. This can be achieved in two ways. The obvious one is to use highly-polarized right-handed electron beams. At high energies the electrons from the W decay are peaked in the backward direction inside the cone $\theta \sim m_W/\sqrt{s}$ (\sqrt{s} is the center of mass energy), so that the W effects can also be significantly cut off by imposing angular cuts in the backward direction of the outgoing electron. It is also expected that the photon beam obtained by back-scattering laser light off the electron beam will not be monochromatic in practice. To simplify our analysis and focus on physics issues, we only consider the case with right-handed polarized electrons and monochromatic high energy photon beams. Another advantage of using e_R beams lies in the fact that the coupling of a right-handed electron/selectron to a chargino is suppressed by the electron mass, and this greatly simplifies the analysis. Furthermore, in $e_R\gamma$ collisions the SM backgrounds to the SUSY signatures most frequently involve a neutrino pair from Z decay carrying away the missing energy, and this can be efficiently eliminated by cuts on the invariant mass of the missing energy around m_Z . Contamination from e_L is expected in actual experiments and can be easily incorporated. Similarly, the folding in of the photon spectrum can be done once its actual shape is known.

II. NEUTRALINO LSP SCENARIO

Within the minimal supersymmetric standard model, the lightest neutralino state χ_1 is often assumed to be the LSP. The mass of χ_1 is generally expected to be at least tens of GeV [12]. We refer to the MSSM as the model with minimal particle content, with R -parity conservation, and without assuming gauge coupling, gaugino mass, or scalar mass unification

at high energy scales [3,4]. In general, the states of the neutralino weak interaction basis mix among themselves and can be transformed into mass eigenstates (χ_i , $i = 1, 2, 3, 4$) by a unitary matrix N' which is defined by the following relation among two-component spinors

$$\chi_i = N'_{ij}\psi'_j, \quad (1)$$

where ψ'_j contains the four states in the weak interaction basis with $j = 1, 2, 3, 4$ referring to the photino ($\tilde{\gamma}$), zino (\tilde{Z}) and two orthogonal linear combinations of the two higgsinos, respectively [4]. Both the masses of the neutralinos and the matrix N' are functions of the Higgs mixing parameter μ , the ratio of the two Higgs' VEVs $\tan\beta$, and the $U(1)_Y$ and $SU(2)_L$ gaugino masses M_1 and M_2 .

A. The Step Function Behavior of the Outgoing Electron Energy Spectrum

We start by considering the simplest case [5,6] where

$$e_R\gamma \rightarrow \tilde{e}_R\chi_1 \quad \text{and} \quad \tilde{e}_R \rightarrow e\chi_1, \quad (2)$$

assuming that χ_1 is gaugino-like. The experimental signature in this case is $e + \cancel{E}$. Throughout this letter we will use the narrow width approximation which is justified because the sparticles decay weakly. In this approximation the differential cross sections can be easily obtained. As the \tilde{e}_R decay is isotropic in its rest frame, boosting back to the lab frame (i.e. the $e\gamma$ CM frame) gives a flat distribution in the electron energy spectrum¹. The standard model background $e\gamma \rightarrow eZ$ with $Z \rightarrow \nu\bar{\nu}$ has an electron energy distribution with a Breit-Wigner shape, and can be eliminated by cuts on the outgoing electron energy.

¹Folding in the photon energy spectrum will in general smear out the step function. However there are still an upper and a lower bound on the outgoing electron energy which can be used to determine $m_{\tilde{e}}$ and m_{χ_1} [6].

The endpoints (E_{\pm}) in the electron energy spectrum can be obtained from kinematical considerations. Inverting these relations gives the masses of the \tilde{e}_R and χ_1 ,

$$m_{\tilde{e}} = \frac{s\sqrt{E_+E_-}}{\sqrt{s}(E_+ + E_-) - 2E_+E_-} \quad (3)$$

$$m_{\chi_1} = \sqrt{m_{\tilde{e}}^2 - 2m_{\tilde{e}}\sqrt{E_+E_-}}. \quad (4)$$

The cross section for the above mentioned selectron production and decay is given by

$$\sigma(e_R^-\gamma \rightarrow e^-\chi_1\chi_1) = |N'_{11} - N'_{12}\tan\theta_W|^4 \sigma(e_R^-\gamma \rightarrow e^-\tilde{\gamma}\tilde{\gamma}), \quad (5)$$

where θ_W is the weak mixing angle. The cross section on the right hand side is obtained in the limit when the photino is the LSP. Therefore, a measurement of the electron energy spectrum (and the total cross section) can determine $m_{\tilde{e}}$, m_{χ_1} and $|N'_{11} - N'_{12}\tan\theta_W|$. Since m_{χ_1} and $N'_{11} - N'_{12}\tan\theta_W$ are functions of μ , $\tan\beta$, M_1 and M_2 only, nontrivial constraints on these four parameters can be obtained.

If χ_1 is higgsino-like and χ_2 is gaugino-like, the reaction may proceed through the gaugino component of χ_2 , $e_R\gamma \rightarrow \tilde{e}_R\chi_2$, with subsequent decays of the selectron ($\tilde{e}_R \rightarrow e\chi_2$, where χ_2 can be real or virtual²) and χ_2 . The electron from \tilde{e}_R decay again has a flat distribution whose endpoints determine $m_{\tilde{e}}$ and m_{χ_2} . The experimental signature depends on the decay modes of χ_2 . For example, the radiative decay $\chi_2 \rightarrow \chi_1\gamma$ can be the dominant mode if χ_1 is higgsino-like and χ_2 is gaugino-like, and the experimental signature is then $e\gamma\gamma + \cancel{E}$. This signature also arises when the gravitino is the LSP and it will be discussed in detail later in Section III. Cuts on the invariant mass of the \cancel{E} around m_Z will get rid of the SM background.

²Note that there is no coupling of the right-handed selectron to charginos in the massless electron limit.

B. Double Step Function Behavior of the Outgoing Electron Energy Spectrum

If both χ_1 and χ_2 contain non-negligible gaugino components, their couplings to the electron and selectron will not be suppressed. Consider first the case in which χ_2 is above the production threshold but is lighter than the \tilde{e}_R , and $\chi_{3,4}$ are heavier than the \tilde{e}_R or are higgsino-like. Then we have

$$e_R\gamma \rightarrow \tilde{e}_R\chi_1 \quad \text{and} \quad \tilde{e}_R \rightarrow e\chi_1/e\chi_2. \quad (6)$$

The electron energy spectrum is now the superposition of two step functions with their endpoints in the $e\gamma$ CM frame E_{\pm}^i ($i = 1, 2$) given by

$$E_{\pm}^i = \frac{\sqrt{s}}{4} \left(1 - \frac{m_{\chi_i}^2}{m_e^2} \right) \left(1 + \frac{m_e^2 - m_{\chi_1}^2}{s} \right) (1 \pm \bar{\beta}_e), \quad (7)$$

where $\bar{\beta}_e$ is the velocity of the \tilde{e}_R , $\bar{\beta}_e = \lambda^{\frac{1}{2}}(s, m_e^2, m_{\chi_1}^2)/(s + m_e^2 - m_{\chi_1}^2)$ and $\lambda(x, y, z) \equiv x^2 + y^2 + z^2 - 2xy - 2yz - 2xz$.

The cross section for the process in Eq. (6) is

$$\begin{aligned} \sigma(e_R\gamma \rightarrow eX) &= \sigma(e_R\gamma \rightarrow \tilde{e}_R\chi_1)[\text{BR}(\tilde{e}_R \rightarrow e\chi_1) + \text{BR}(\tilde{e}_R \rightarrow e\chi_2)] \\ &= \sigma(e_R\gamma \rightarrow \tilde{e}_R\tilde{\gamma}_1)|N'_{11} - N'_{12}\tan\theta_W|^2 \times [\text{BR}(\tilde{e}_R \rightarrow e\chi_1) + \text{BR}(\tilde{e}_R \rightarrow e\chi_2)] \end{aligned} \quad (8)$$

where $X = \chi_1\chi_1$, $\chi_1\chi_2$, and $\tilde{\gamma}_1$ denotes a photino with mass m_{χ_1} . The branching ratios are given by

$$\text{BR}(\tilde{e}_R \rightarrow e\chi_i) = \frac{|N'_{i1} - N'_{i2}\tan\theta_W|^2 (1 - m_{\chi_i}^2/m_e^2)^2}{\sum_{j=1,2} |N'_{j1} - N'_{j2}\tan\theta_W|^2 (1 - m_{\chi_j}^2/m_e^2)^2} \quad i = 1, 2. \quad (9)$$

The electron energy spectrum is plotted in Fig. 1 for some representative masses.

Besides the signature $e + \cancel{E}$ from the chain $e_R\gamma \rightarrow \tilde{e}_R\chi_1 \rightarrow e\chi_1\chi_1$, the experimental signature also depends on the χ_2 decay modes from the chain $e_R\gamma \rightarrow \tilde{e}_R\chi_1 \rightarrow e\chi_1\chi_2$. The signatures that we will focus on are $eX + \cancel{E}$ where X contains no electrons and thus measurement of the single electron energy spectrum can be easily performed. These include

for example, the following χ_2 decay modes: i) the invisible decay $\chi_2 \rightarrow \chi_1 \nu \bar{\nu}$ that gives rise to the signature $e + \cancel{E}$; ii) the radiative decay $\chi_2 \rightarrow \chi_1 \gamma$ which leads to the signature $e \gamma + \cancel{E}$; and iii) the decay through the lightest Higgs boson, $\chi_2 \rightarrow \chi_1 h \rightarrow \chi_1 b \bar{b}$, with the signature $e j_b j_b + \cancel{E}$, where j_b denotes a b -jet. If $m_{\chi_2} > m_h$ then the b -quarks will reconstruct the Higgs. A comprehensive study of the various decay modes of χ_2 can be found in Ref. [13]. In the first two cases, the SM backgrounds involve escaped neutrino pairs from Z decay and can be effectively eliminated by imposing cuts on the invariant mass of the \cancel{E} around m_Z . In the third case the \cancel{E} and/or the two b -jets should reconstruct the Z for the SM background.

The masses of the selectrons and neutralinos, $m_{\tilde{e}}$, m_{χ_1} and m_{χ_2} , can be determined once the endpoints of the two step functions E_{\pm}^i ($i = 1, 2$) are measured from the electron energy spectrum (see Fig. 1 for illustration). Furthermore, measurement of the partial cross section for $e_R \gamma \rightarrow e \chi_1 \chi_1$ from one of the two step functions fixes a relation between $|N'_{11} - N'_{12} \tan \theta_W|$ and $|N'_{21} - N'_{22} \tan \theta_W|$. Measurement of the other step function gives a lower bound on the partial cross section for $e_R \gamma \rightarrow e \chi_1 \chi_2$ and thus lower bounds on both $|N'_{11} - N'_{12} \tan \theta_W|$ and $|N'_{21} - N'_{22} \tan \theta_W|$. If χ_2 decays only via the above three modes, then both $|N'_{11} - N'_{12} \tan \theta_W|$ and $|N'_{21} - N'_{22} \tan \theta_W|$ can be determined. Recall that m_{χ_1} , m_{χ_2} and N'_{ij} are functions of μ , $\tan \beta$, M_1 and M_2 . These latter four parameters can thus be solved from the four known quantities, m_{χ_1} , m_{χ_2} , $|N'_{11} - N'_{12} \tan \theta_W|$ and $|N'_{21} - N'_{22} \tan \theta_W|$.

The other possibilities, including the case $e_R \gamma \rightarrow \tilde{e}_R \chi_1 / \chi_2$ with $\tilde{e}_R \rightarrow e \chi_1 / \chi_2$ where the outgoing electron energy spectrum is a superposition of four step functions, can be similarly analyzed. The experimental signatures are in general more involved, and no further information about the MSSM can be gained beyond $m_{\tilde{e}}$, μ , $\tan \beta$, M_1 and M_2 .

III. GRAVITINO AS LSP

Models of low energy gauge-mediated supersymmetry breaking [7] have quite a different mass spectrum from the MSSM. In particular, it is very likely that this SUSY breaking

scheme would lead to the gravitino being the LSP with the NLSP being either the lightest neutralino or a right-handed slepton, which we will assume to be \tilde{e}_R . The gravitino gets its mass via the super Higgs mechanism (assuming zero cosmological constant) with $m_{\tilde{G}} = \kappa d / \sqrt{6} \simeq 1.7 (\sqrt{d}/100 \text{ TeV})^2 \text{ eV}$, where d denotes the scale of supersymmetry breaking and $\kappa = \sqrt{8\pi G_{\text{Newton}}}$. In models of low energy dynamical supersymmetry breaking (DSB), \sqrt{d} ranges from 100 TeV to a few thousand TeV, and the gravitino mass takes its values in the eV to keV range.

It is well known [8] that a light gravitino can couple to matter with weak interaction strength instead of gravitationally via its Goldstino component. The Goldstino has couplings to a particle and its superpartner determined by the supersymmetric analog of the Goldberger-Treiman relation, and is proportional to $\Delta m / 2d$, where Δm measures the mass splitting between an ordinary particle and its superpartner. This allows the NLSP (or any heavier sparticle) to decay into its SM partner and a gravitino with universal coupling strength.

Consider first the case of the lightest neutralino χ_1 being the NLSP. Since the cross section for direct production of a higgsino-like χ_1 via $e\gamma$ collision is negligible due to the tiny Yukawa coupling of the electron, we will restrict our analysis to a gaugino-like χ_1 and focus on the dominant process,

$$e_R \gamma \rightarrow \tilde{e}_R \chi_1 \rightarrow e \chi_1 \chi_1 \quad \text{and} \quad \chi_1 \rightarrow \gamma \tilde{G}, \quad (10)$$

where \tilde{G} denotes the gravitino. If $m_{\chi_1} > m_Z$, the decay $\chi_1 \rightarrow Z \tilde{G}$ can proceed through the zino component of χ_1 ; For most of the parameter space, the photonic decay will be the dominant mode, with its width given by [8]

$$\Gamma(\chi_1 \rightarrow \gamma \tilde{G}) = |N'_{11}|^2 m_{\chi_1}^5 / (8\pi d^2) \simeq 1.1 \times 10^{-2} |N'_{11}|^2 (m_{\chi_1}/100 \text{ GeV})^5 (1 \text{ eV}/m_{\tilde{G}})^2 \text{ eV}, \quad (11)$$

where N'_{11} is defined in Eq. (1). This translates into a decay distance for χ_1 given by

$$D(\chi_1 \rightarrow \gamma \tilde{G}) \simeq 1.8 \times 10^{-3} |N'_{11}|^{-2} \sqrt{E_{\chi_1}^2/m_{\chi_1}^2 - 1} (m_{\tilde{G}}/1 \text{ eV})^2 (100 \text{ GeV}/m_{\chi_1})^5 \text{ cm}, \quad (12)$$

where E_{χ_1} is the energy of χ_1 . For $m_{\tilde{G}} \sim \text{eV} - \text{keV}$ and $m_{\chi_1} \sim 100\text{GeV}$, the decay length of χ_1 ranges from hundreds of microns to tens of meters [9]. It is therefore possible to observe the background-free signature “ $e + \text{displaced } \gamma\gamma + \cancel{E}$ ”. The electron energy spectrum has the characteristic flat distribution. If the supersymmetry breaking scale \sqrt{d} is several thousand TeV then the decay will occur outside the detector. In this case the signature would be $e + \cancel{E}$ and would be indistinguishable from the neutralino scenario considered in Section II.

If the NLSP decays quickly, the experimental signature will be $e\gamma\gamma + \cancel{E}$ and has the SM background from $e_R\gamma \rightarrow e\gamma\gamma Z$ with $Z \rightarrow \nu\bar{\nu}$. Cuts on the invariant mass of the \cancel{E} at m_Z should again allow one to remove the SM background.

The signature $e\gamma\gamma + \cancel{E}$ can also arise in the MSSM where $e_R\gamma \rightarrow \tilde{e}_R\chi_2$ with $\tilde{e}_R \rightarrow e\chi_2$ and $\chi_2 \rightarrow \chi_1\gamma$. This happens if χ_1 is higgsino-like, χ_2 is gaugino-like, and χ_2 has a large branching ratio into χ_1 and γ . Examination of the neutralino mass matrix [13,10] shows this could indeed occur for certain ranges of parameter space (for example, $\tan\beta \simeq 1$ and $-\mu = |\mu| < M_1 \simeq M_2$, as noted in ref. [10]). The electron energy spectrum is in both scenarios a step function whose endpoints allow for a determination of $m_{\tilde{e}}$ and m_{χ_1} (m_{χ_2}) in DSB (MSSM) (cf. Eqs. (3) and (4)). The endpoints in the photon energy spectrum can then be used to determine the LSP mass and therefore to distinguish unambiguously between the DSB and MSSM scenarios. We now derive the photon energy distributions in both scenarios, denoting the LSP and NLSP by X_1 and X_2 for convenience of presentation.

Consider first the reaction chain $e_R\gamma \rightarrow \tilde{e}_RX_2$ with $X_2 \rightarrow X_1\gamma$ ³. The photon distribution in the rest frame of the X_2 is isotropic. Boosting back to the lab frame ($e\gamma$ CM frame) gives a flat energy distribution for the photon,

³We are assuming for simplicity that the branching ratios for $X_2 \rightarrow X_1\gamma$ and $\tilde{e}_R \rightarrow eX_2$ are both equal to one. The endpoints of the electron and photon energy spectra, which are used to distinguish between the MSSM and DSB scenarios, are independent of this assumption.

$$\frac{d\sigma_1}{dE} = \frac{\sigma(e_R\gamma \rightarrow \tilde{e}_R X_2)}{E_1^+ - E_1^-} \quad \text{for } E_1^- \leq E \leq E_1^+, \quad (13)$$

where E_1^+ and E_1^- are the maximum and minimum photon energies in the lab frame,

$$E_1^\pm = \frac{\sqrt{s}}{4} \left(1 - \frac{m_{X_1}^2}{m_{X_2}^2}\right) \left(1 - \frac{m_e^2 - m_{X_2}^2}{s}\right) (1 \pm \beta_{X_2}), \quad (14)$$

and where $\beta_{X_2} = \lambda^{\frac{1}{2}}(s, m_e^2, m_{X_2}^2)/(s - m_e^2 + m_{X_2}^2)$ is the velocity of X_2 in the lab frame.

We now turn to the second reaction chain, $e_R\gamma \rightarrow \tilde{e}_R X_2$ with the subsequent decays $\tilde{e}_R \rightarrow e X_2$ and $X_2 \rightarrow X_1 \gamma$. The photon energy distribution in the lab frame can be obtained by boosting back first to the \tilde{e}_R rest frame from the X_2 rest frame followed by a boost into the lab frame, and it is given by

$$\frac{d\sigma_2}{dE} = \frac{\sigma(e_R\gamma \rightarrow \tilde{e}_R X_2)}{E_2^- + E_2^+ - E_2^\alpha - E_2^\beta} \times \begin{cases} \ln E/E_2^- & \text{if } E_2^- \leq E \leq E_2^\alpha \\ \ln \frac{1+\beta}{1-\beta} & \text{if } E_2^\alpha \leq E \leq E_2^\beta \\ \ln E_2^+/E & \text{if } E_2^\beta \leq E \leq E_2^+ \\ 0 & \text{otherwise} \end{cases} \quad (15)$$

where $E_2^\alpha = \min(E_2^A, E_2^B)$, $E_2^\beta = \max(E_2^A, E_2^B)$, and

$$E_2^{\pm, A, B} = \frac{\sqrt{s}}{8} \left(1 + \frac{m_{X_2}^2}{m_e^2}\right) \left(1 - \frac{m_{X_1}^2}{m_{X_2}^2}\right) \left(1 + \frac{m_e^2 - m_{X_2}^2}{s}\right) (1 \pm \beta'_{X_2}) (1 \pm \beta_e), \quad (16)$$

with E_2^- , E_2^A , E_2^B and E_2^+ given by the $--$, $-+$, $+-$ and $++$ combinations respectively; and $\beta = \min(\beta'_{X_2}, \beta_e)$, where $\beta'_{X_2} = (m_e^2 - m_{X_2}^2)/(m_e^2 + m_{X_2}^2)$ is the X_2 velocity measured in the rest frame of the \tilde{e}_R and $\beta_e = \lambda^{\frac{1}{2}}(s, m_e^2, m_{X_2}^2)/(s + m_e^2 - m_{X_2}^2)$ is the \tilde{e}_R velocity in the lab frame.

The single photon energy spectrum measured in a real experiment has contributions from both photons in the above discussed reaction chains. Since there is no correlation in the energies of the two decay photons, the single photon energy spectrum is simply given by the average $d\sigma/dE_\gamma = (d\sigma_1/dE + d\sigma_2/dE)/2$. Fig. 2 gives the single photon energy spectra for both DSB and MSSM where the difference between the two curves is due to the LSP mass.

Alternatively, the spectrum of the sum of the two photon energies can be measured and used to distinguish between the MSSM and DSB scenarios. The differential cross section can be derived as

$$\frac{d\sigma}{d(E_1^\gamma + E_2^\gamma)} = \frac{1}{E_1^+ - E_1^-} \int_{E-E_1^+}^{E-E_1^-} H(x) dx \quad \text{for } E_1^- + E_2^- \leq E \leq E_1^+ + E_2^+, \quad (17)$$

where $H(E) \equiv d\sigma_2/dE$ as given by Eq. (15). These are plotted for the MSSM and DSB scenarios in Fig. 3, where their difference is essentially in the endpoints of the sum of the two photon energies as a result of different LSP masses.

The analysis for $\chi_1 \rightarrow Z\tilde{G}$ in DSB is similar to the discussion presented above for the photonic decay of χ_1 and will not be repeated here.

The second possibility in low energy DSB models is that the right-handed slepton plays the role of NLSP. We will assume this to be a \tilde{e}_R . The decay chain of interest is $e_R\gamma \rightarrow \tilde{e}_R\chi_1$ with $\chi_1 \rightarrow \tilde{e}_R e^+/\tilde{e}_R^* e^-$ and $\tilde{e}_R(\tilde{e}_R^*) \rightarrow e^-(e^+)\tilde{G}$. The background-free signatures would be either three charged tracks without \cancel{E} if the selectrons decay outside the detector, or “ $e^+ + \text{displaced } e^-e^- + \cancel{E}$ ” and “ $e^- + \text{displaced } e^+e^- + \cancel{E}$ ” if the selectrons decay inside the detector. For short decay lengths of the selectron, the signature will be $e^+e^-e^- + \cancel{E}$. The SM background in this case is quite involved, especially due to the contributions from W ’s, and in general it cannot be eliminated by imposing cuts on the invariant mass of \cancel{E} . It appears that it may in fact be difficult to observe this signal over the background.

IV. CONCLUSION

Within the frameworks of the minimal supersymmetric standard model and low energy dynamical supersymmetry breaking, we have explored the unique signatures arising from high energy $e_R\gamma$ collisions, $e + \cancel{E}$ and $eX + \cancel{E}$ ($X = \gamma, \gamma\gamma, j_b j_b, e^+e^-, \dots$). In cases where there is only one charged particle (the electron) in the final state, the electron energy spectrum enjoys a flat distribution or superposition of flat distributions.

Using a right-handed electron beam not only serves to eliminate the large W background, but also offers an efficient way to separate out the remaining SM backgrounds by cuts on the invariant mass of the \cancel{E} carried away by a neutrino pair from Z decay. The analysis is further simplified by the fact that only right-handed selectrons will be produced and that these cannot decay into charginos.

In the MSSM, where the lightest neutralino χ_1 is assumed to be the LSP, we have observed that a double step function behavior of the electron energy spectrum may allow for a complete determination of the selectron mass and four basic parameters of the MSSM μ , $\tan\beta$, M_1 and M_2 , and therefore provides an independent check of relations among these parameters derived from higher scale physics like GUT. In low energy DSB, where the gravitino is the LSP, the NLSP can be either the lightest neutralino or a right-handed slepton. Depending on the supersymmetry breaking scale, the NLSP could have long decay lengths and therefore could give rise to background-free signatures, $e +$ displaced $\gamma\gamma + \cancel{E}$ for the former, and three charged tracks or $e^\pm +$ displaced $e^\mp e^- + \cancel{E}$ for the latter. Even when the dynamics is such that background-free decays do not occur, we find that the outgoing electron and photon energy spectra can be very useful diagnostic tools for distinguishing between the DSB and MSSM scenarios.

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Figure Captions

Figure 1: (a) The “double step function” for the outgoing electron energy spectrum in the MSSM (cf. Eq. (6)). The kinematics is such that χ_2 cannot be produced initially, but may be produced from the decay of the selectron. For the purpose of illustration we have taken the LSP and NLSP to be the photino and zino, with masses 150 and 200 GeV, respectively, and have set $\sqrt{s} = 500$ GeV and $m_{\tilde{e}} = 300$ GeV. Under these assumptions the selectron can only decay into the photino and zino modes, with branching ratios 86% and 14%, respectively. The solid line shows the sum of the two step functions and the dotted and dashed lines show the individual step functions due to photino and zino production. (b) The double step function for the same set of parameters as in (a) except that now $m_{\chi_2} = 250$ GeV. In this case the branching ratios for the photino and zino modes are 95% and 5%.

Figure 2: The single photon energy spectrum for $e_R\gamma \rightarrow \tilde{e}_R X_2 \rightarrow e X_2 X_2 \rightarrow e X_1 X_1 \gamma\gamma$ where X_1 is the LSP and X_2 is the NLSP. In this plot we have taken X_2 to be the photino and we have set $\sqrt{s} = 500$ GeV, $m_{\tilde{e}} = 200$ GeV and $m_{X_2} = 100$ GeV. The solid line corresponds to an LSP mass of 50 GeV (in the MSSM scenario) and the dashed line corresponds to a massless LSP (which is a good approximation in the gravitino scenario.) If X_2 is not pure photino then the solid (dashed) curve needs simply to be scaled by an overall factor of $|N'_{21} - N'_{22}\tan\theta_W|^2$ ($|N'_{11} - N'_{12}\tan\theta_W|^2$).

Figure 3: The differential spectrum for the sum of the two photon energies, $d\sigma/d(E_1^\gamma + E_2^\gamma)$, for the reaction $e_R\gamma \rightarrow \tilde{e}_R X_2 \rightarrow e X_2 X_2 \rightarrow e X_1 X_1 \gamma\gamma$, using the same parameters as in Fig. 2.

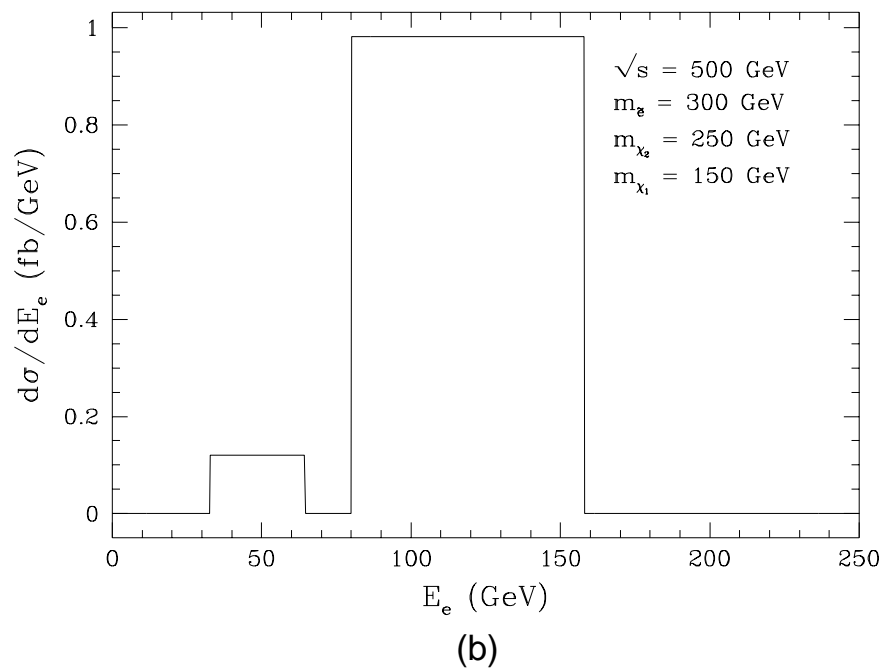
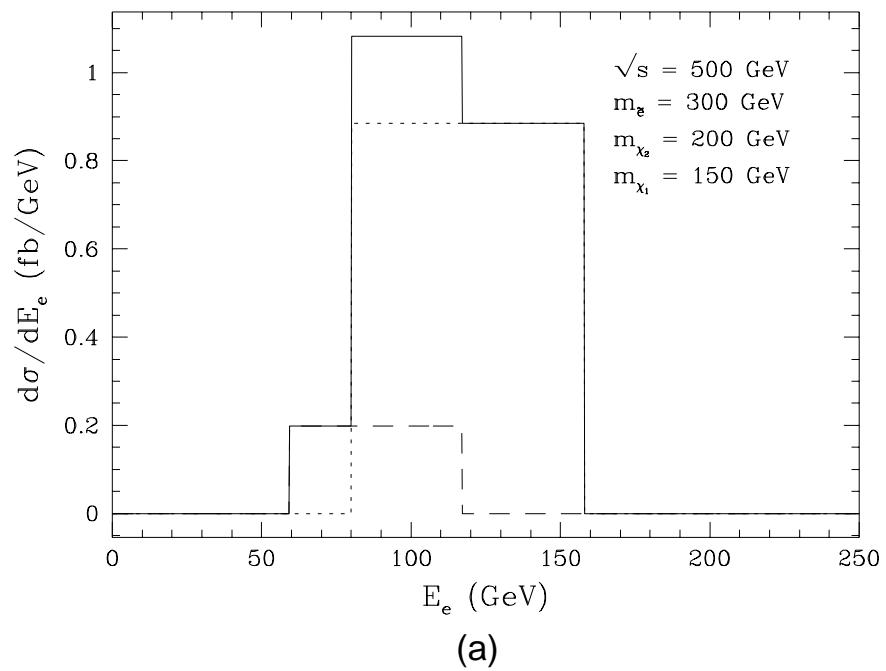


Figure 1

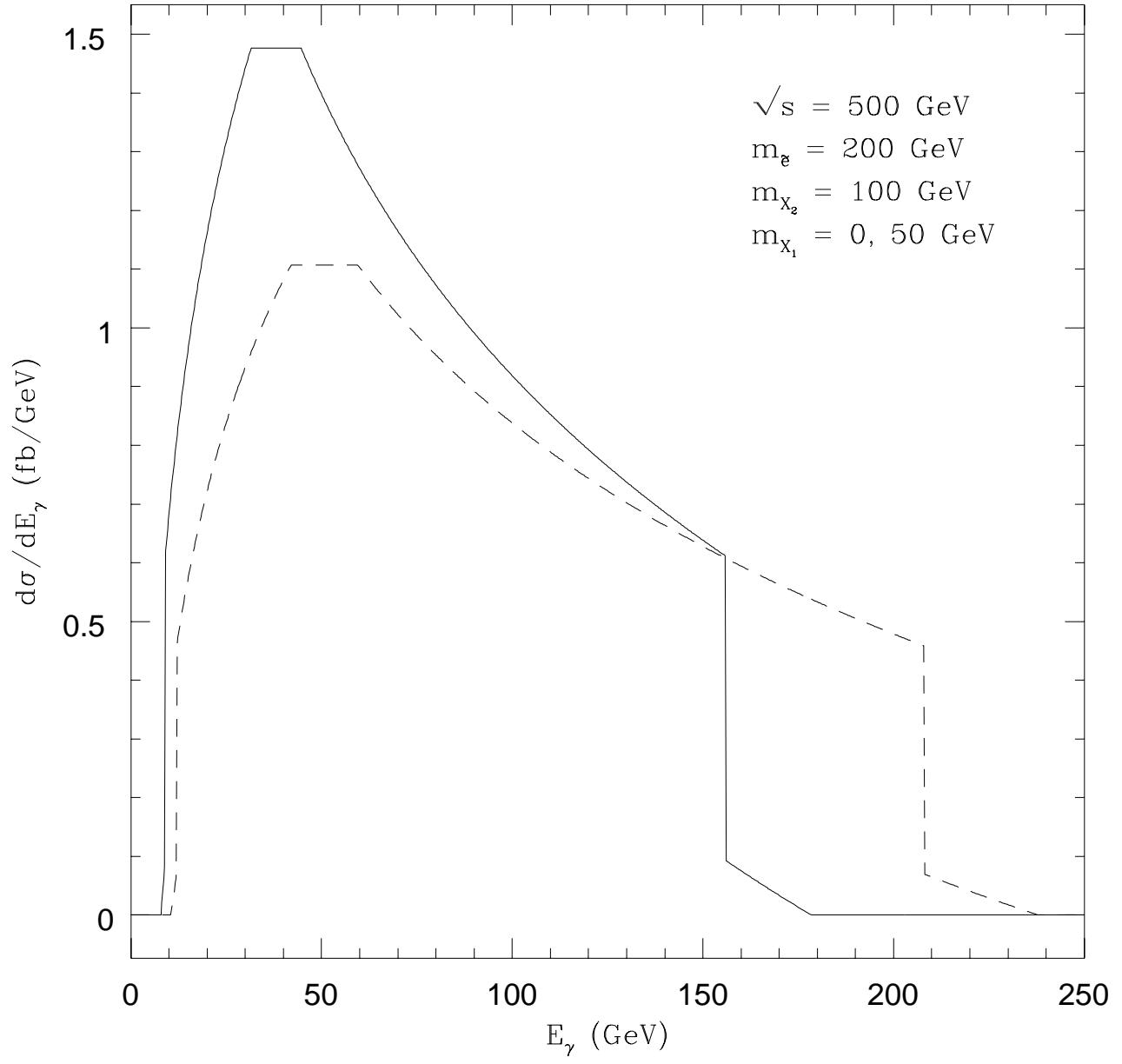


Figure 2

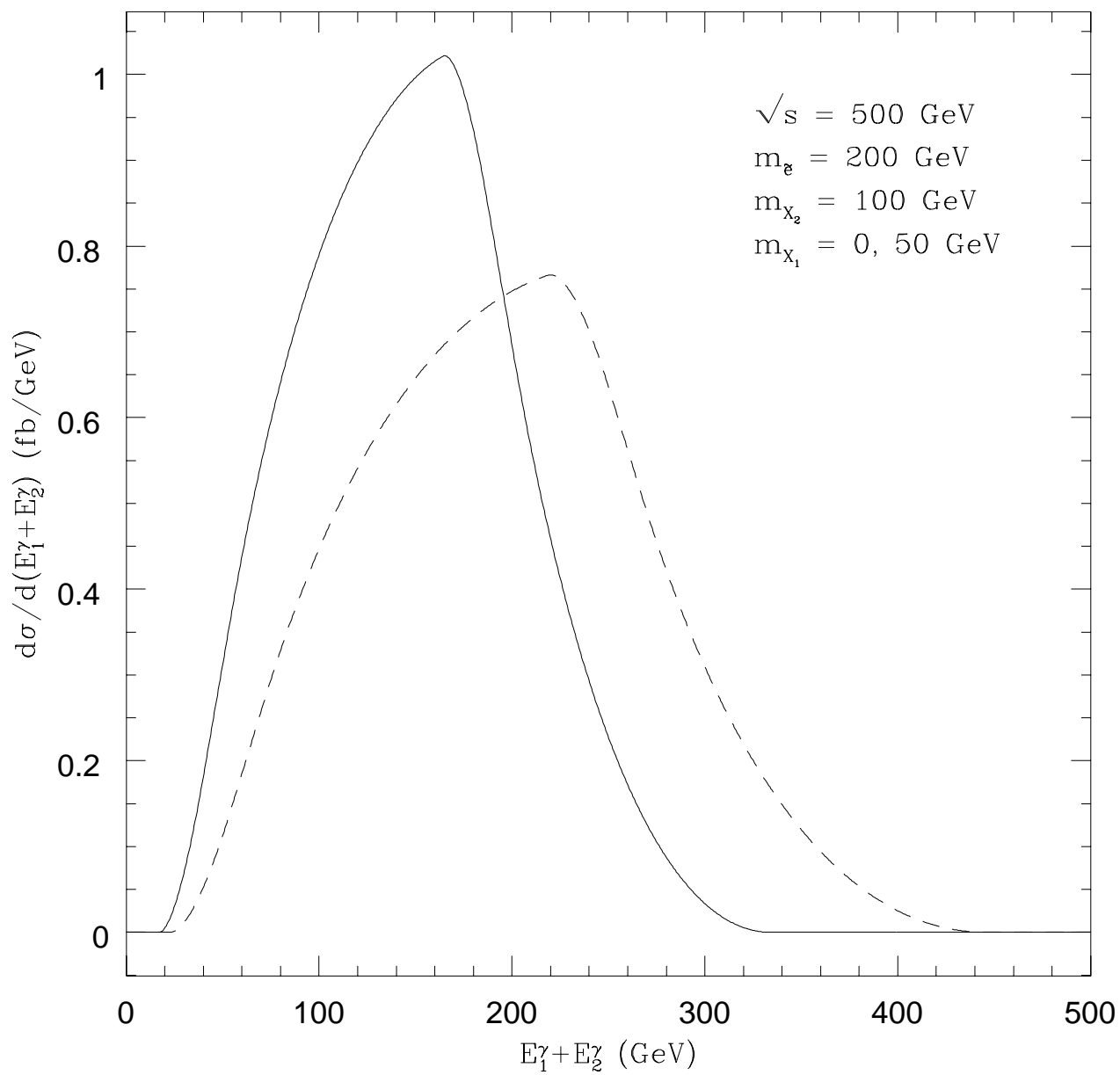


Figure 3